

WU-HEP-12-08
November, 2012

Instant uplifted inflation

:A solution for a tension between inflation and SUSY breaking scale

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Abstract

The Hubble parameter during the inflationary era must be smaller than the gravitino mass if the moduli are stabilized by the Kachru-Kallosh-Linde-Trivedi mechanism. This condition represents the difficulty to combine the low scale SUSY breaking and the high scale inflation. We propose a simple mechanism which can naturally separate the inflation scale from the SUSY breaking scale today.

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1 Introduction

Cosmological inflation [1] can solve some problems of standard cosmology (e.g. flatness and horizon problem), so it is strongly favored theoretically. On the other hand, typical slow-roll inflation models provide source of the primordial density perturbation which is almost scale invariant (for recent review see Ref. [2]). Such the density perturbation is favored for the consistency with the CMB (cosmic microwave background) observation. Therefore the observation also favors the existence of an inflationary era.

In particle physics, there are also some theoretical problems such as the gauge hierarchy problem in the standard model (SM) of elementary particles . Supersymmetry (SUSY) is the promising solution for the gauge hierarchy problem due to the absence of quadratic divergences. The observational and experimental data indicate that SUSY must be broken, and the SUSY breaking scale should be above the electroweak scale. In the minimal supersymmetric standard model (MSSM), a mass of the Z boson is related to a soft SUSY breaking mass of the Higgs field. Recently, Higgs-like boson was discovered at the large hadron collider (LHC) with its mass $m_H \sim 125\text{GeV}$. Then, if we consider SUSY as a solution for the gauge hierarchy problem, a TeV-scale SUSY breaking model is favored by the naturalness. In addition, MSSM with the low scale SUSY breaking provides good candidates for the dark matter. Therefore, models with the low scale SUSY breaking are fascinating.

In order to discuss both cosmological inflation and the SM together, we have to use the self-consistent quantum gravity theory. Superstring theory is the most promising candidate for the quantum gravity theory which has a possibility to explain from the cosmological observation to the high energy experiments. Superstring theory predicts the six dimensional extra space, whose volume and shape are determined by vacuum expectation values (VEVs) of the moduli fields. In the four dimensional (4D) effective theories, parameters in the SM such as gauge coupling constants are also determined by those VEVs, and then the moduli stabilization is an important issue.

The Kachru-Kallosh-Linde-Trivedi (KKLT) model [3] is a well known moduli stabilization scenario in the superstring models, where the stabilization mechanism consists of three steps. First, it is assumed that the dilaton and complex structure moduli are stabilized through three form fluxes at a high scale. At the second step, a Kähler moduli dependent term in the superpotential is introduced assuming a certain non-perturbative effect. Thus, Kähler moduli is stabilized by such superpotential. But, the minimum of the scalar potential is negative valued. At the third step, the anti de Sitter (AdS) vacuum is uplifted to the de Sitter vacuum by a SUSY breaking term with a positive energy.

It was pointed out in Ref. [4] that in KKLT type models the Hubble parameter H and the gravitino mass $m_{3/2}$ must satisfy a condition $H \lesssim |m_{3/2}|$ to stabilize the moduli during the inflationary era. In that case, H must be below the TeV scale to construct a low scale SUSY model such as $m_{3/2} \sim \mathcal{O}(1)\text{TeV}$. Then, it is difficult to generate the scalar perturbation consistent with the observation. This problem was pointed out by Kallosh and Linde which we call the Kallosh-Linde (KL) problem.

The KL problem occurs if moduli are not inflatons. Independently, from the minimalistic point of view, it seems natural to consider the case that the moduli play a roll of inflaton, and such models were suggested. (For example, see Refs. [5], [6], and [7].) They are successful models from the viewpoint of the inflation, however do not realize a low scale SUSY breaking. In Refs. [8] and [9], it was shown that the modulus can not have the inflationary de Sitter point without SUSY breaking terms if its Kähler potential is given by $K = -n \log(T + \bar{T})$ for $0 < n \leq 3$. (We call such moduli fields as the volume type moduli.) Then, in typical moduli inflation models, we need to add the SUSY breaking terms to realize the inflationary de Sitter point. Therefore the SUSY breaking scale is related to the Hubble scale in the same way as the KL problem. In addition, the moduli have a possibility to overshoot the minimum and to be destabilized after inflation. This is so-called the overshooting problem. This problem also causes a difficulty in moduli inflation models.

In this paper, we propose a new mechanism to combine high scale inflation with low scale SUSY breaking. To achieve this goal, there are two important ingredients. One is a SUSY breaking field Y which has a superpotential term $W = \mu_Y^2 Y e^{-c_Y T}$. That kind of terms generates the F-term varying exponentially in terms of T . So, even if such a F-term makes a inflationary de Sitter point at the high scale, the F-term becomes small enough as the modulus rolling into a mildly large VEV. The other one is the non-perturbative

superpotential with positive exponents [11]. As pointed out in Ref. [11], it can prevent the overshooting problem without fine-tuned initial conditions and parameters.

This paper is organized as follows. In Sec. 2, we review the reason why it is difficult to combine high scale inflation models with the low energy SUSY breaking. Then, we discuss the model with the two ingredients mentioned above and find that those can separate the SUSY breaking scale from the inflation scale in Sec. 3. In Sec. 4, we show explicit models with different types of moduli stabilization. Finally, we conclude in Sec. 5.

2 Review of the KL problem

In the typical inflation models with moduli, the gravitino mass $m_{3/2}$ today are related to the Hubble parameter during the inflationary era. Therefore, a low scale SUSY breaking model leads to a low scale inflation model. There are two situations. (a) One is that the moduli are not inflatons. (b) The other one is that the moduli are inflaton.

First, we review the situation (a) (so-called the KL problem). To discuss concretely, we consider a simple KKLT model [3] that a modulus field $T = \sigma + i\alpha$ has the Kähler potential and superpotential as follows:

$$K = -3 \log(T + \bar{T}), \quad (2.1)$$

$$W = w_0 + A e^{-aT}. \quad (2.2)$$

The F-term scalar potential in 4D supergravity (SUGRA) can be written in terms of W and K in the following form:

$$V = e^K (D_I W K^{I\bar{J}} D_{\bar{J}} \bar{W} - 3|W|^2), \quad (2.3)$$

where $D_I W = \partial_I W + \partial_I K W$, the indices I, \bar{J} denote the corresponding chiral superfields Q^I and their conjugates $\bar{Q}^{\bar{J}}$ respectively, and ∂_I represents a derivative with respect to a lowest component of a chiral superfield Q^I .

In this case, the SUSY condition $D_T W = 0$ is satisfied at the minimum. Then, the VEV of the scalar potential at that point is given by (in the Planck unit $M_{\text{pl}} = 1$)

$$\langle V \rangle_{\text{AdS}} = -3e^{\langle K \rangle} |\langle W \rangle|^2 = -3m_{3/2}^2. \quad (2.4)$$

To vanish the VEV of the scalar potential at the minimum (to be precisely, the value has to be $\Lambda = \mathcal{O}(10^{-120})$), we have to add the SUSY breaking terms¹ to uplift the minimum:

$$\begin{aligned} V &= e^K (D_I W K^{I\bar{J}} D_{\bar{J}} \bar{W} - 3|W|^2) + V_{\text{uplift}}, \\ V_{\text{uplift}} &\sim |3m_{3/2}^2|. \end{aligned} \quad (2.5)$$

¹In the original model [3], they added the uplifting term from anti-D3 branes. We can also uplift the potential with the non-vanishing F-terms as Refs. [12].

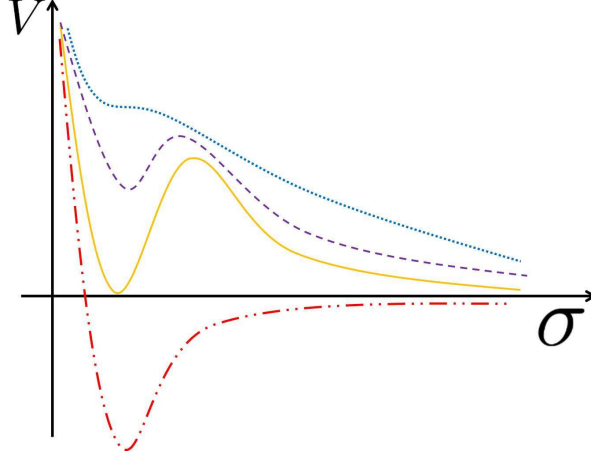


Figure 1: An illustration of the F-term potential for a typical KKLT model drawn by the red dashed-two-dotted line. (See Eq. (2.4).) The uplifted potential is drawn by a yellow normal line. (See Eq. (2.5).) As the inflationary potential becomes larger (see Eq. (2.7)), the barrier becomes smaller or disappears. (Drawn by a purple dashed line and blue dotted line respectively.)

As shown schematically in Fig. 1, the potential has a barrier after uplifting. The height of barrier is approximately given by

$$V_B \sim \mathcal{O}(|\langle V_{\text{AdS}} \rangle|) \sim \mathcal{O}(m_{3/2}^2). \quad (2.6)$$

The scalar potential is generalized to a model that contains an inflaton. For simplicity, we do not consider the case that the modulus and the inflaton ϕ have mixing term in the Kähler potential and the superpotential. Then, the scalar potential becomes a following form:

$$\begin{aligned} V &= e^K (D_\phi W K^{\phi\bar{\phi}} D_{\bar{\phi}} \bar{W} + D_T W K^{T\bar{T}} D_{\bar{T}} \bar{W} - 3|W|^2) + V_{\text{uplift}}, \\ &\sim \frac{1}{8\sigma^3} (D_T W K^{T\bar{T}} D_{\bar{T}} \bar{W} - 3|W(T)|^2) + V_{\text{uplift}} + \frac{V_{\text{inf}}(\phi)}{8\sigma^3} \end{aligned} \quad (2.7)$$

where $V_{\text{inf}}(\phi)$ denotes the inflaton dependent terms and we assume that $V_{\text{inf}}(\phi)$ vanishes at the minimum. Then we find that the term $\frac{V_{\text{inf}}(\phi)}{8\sigma^3}$ has the positive value during the inflationary era. So, that term plays the same role as the uplifting term V_{uplift} . As the inflaton potential becomes larger than $|\langle V \rangle_{\text{AdS}}|$, the scalar potential at the minimum during inflation with respect to the moduli becomes higher and the barrier becomes smaller or disappears. (See Fig. 1.) To avoid such a destabilization, the Hubble parameter needs to satisfy the condition :

$$H \lesssim m_{3/2}. \quad (2.8)$$

This is the KL problem.

Kallosh and Linde pointed out the KL problem in Ref. [4], and suggested a simple solution². It is called the KL model which contains the following alternative superpotential terms:

$$W_{\text{KL}} = w_0 + Ae^{-aT} - Be^{-bT}, \quad (2.9)$$

$$w_0 = B \left(\frac{aA}{bB} \right)^{\frac{b}{b-a}} + A \left(\frac{aA}{bB} \right)^{\frac{a}{b-a}}. \quad (2.10)$$

The stationary point satisfies $D_TW = 0$. (We denote the value of modulus at this point as $T = T_{\text{min}}$.) Unlike the original KKLT model, we find that $W = 0$ at $T = T_{\text{min}}$ because of the condition (2.10). So, the minimum is a Minkowski vacuum, and the height of barrier is not related to the gravitino mass $m_{3/2}$. This model seems simple, however the condition (2.10) requires the fine-tuning of a parameter w_0 ³. Secondly we review the problems in the situation (b). In this case, the general conditions (2.20) shown later are necessary for a realization of the inflationary de Sitter point. That constraint is originated in an inequality (2.19) studied in detail in Refs. [8], [9] and shown later. Based on Ref. [9], we discuss about that condition. In the following discussion, we use the quantities defined by

$$G \equiv K + \log |W|^2, \quad (2.11)$$

$$\gamma \equiv \frac{V}{3e^G} = \frac{V}{3m_{3/2}^2}, \quad (2.12)$$

$$f_I \equiv \frac{G_I}{\sqrt{G^{J\bar{K}}G_JG_{\bar{K}}}}, \quad (2.13)$$

$$G^{I\bar{J}} \equiv (G_{I\bar{J}})^{-1} = (\partial_I \partial_{\bar{J}} G)^{-1}, \quad (2.14)$$

$$R_{I\bar{J}K\bar{L}} \equiv \partial_I \partial_{\bar{J}} G_{K\bar{L}} - G^{M\bar{N}} \partial_{\bar{J}} G_{M\bar{L}} \partial_I G_{K\bar{N}}, \quad (2.15)$$

$$\hat{\sigma}(f^I) \equiv \frac{2}{3} - R_{I\bar{J}K\bar{L}} f^I f^{\bar{J}} f^K f^{\bar{L}}. \quad (2.16)$$

Here we consider the general case in which inflatons do not have a canonical kinetic term.

²Recently, some alternative models to solve the KL problem were suggested by Refs. [13], [14].

³Some of recent works [15], [16] show that the KL model requires the fine-tuning as the same order of tuning as the original KKLT model if the small constant term w_0 is induced by tuning the flux. However, Ref. [17] showed that we can realize the “smallness” of w_0 without fine-tuning if the constant superpotential induced by the flux is zero and a dilaton S dependent gaugino condensation term provides an effective constant term $\langle Ae^{-aS} \rangle$. Then the smallness of w_0 is not a consequence of a fine-tuning.

Generalized slow-roll parameters are given by

$$\epsilon = \frac{\nabla_I V G^{I\bar{J}} \nabla_{\bar{J}} V}{V^2}, \quad (2.17)$$

$$\eta = \text{minimum eigenvalue of } \mathbf{M} \quad (2.18)$$

$$\mathbf{M} = \frac{1}{V} \begin{pmatrix} \nabla^I \nabla_J V & \nabla^I \nabla_{\bar{J}} V \\ \nabla^{\bar{I}} \nabla_J V & \nabla^{\bar{I}} \nabla_{\bar{J}} V \end{pmatrix}$$

where ∇_I is the covariant derivative of the Kähler manifold whose metric is given by $G_{I\bar{J}}$. Because of the fact that η is the minimum eigenvalue of the matrix \mathbf{M} , we find the upper bound on η :⁴

$$\eta \leq -\frac{2}{3} + \frac{4\sqrt{\epsilon}}{\sqrt{3(1+\gamma)}} + \frac{\gamma\epsilon}{1+\gamma} + \frac{1+\gamma}{\gamma} \hat{\sigma}(f^I) \sim \frac{1+\gamma}{\gamma} \hat{\sigma}(f^I) - \frac{2}{3}. \quad (2.19)$$

For achieving successful inflation, $|\eta|$ and ϵ must be small during inflation. As shown in Ref. [9], we find $\hat{\sigma}(f^I) \leq 0$ for $K = -3 \log(T + \bar{T})$. This leads the relation $\eta \lesssim -\frac{2}{3}$ namely $|\eta| \gtrsim \frac{2}{3}$. In Ref. [8], it was pointed out that we see this relation in any case for $K = -n \log(T + \bar{T})$ $0 < n \leq 3$. Therefore, it seems that we can not realize the successful moduli inflation. However, we can avoid this claim if the scalar potential includes F-terms of the other fields [18] or explicit SUSY breaking terms (e.g. anti-D3 brane) [8]. Then, such terms may give the inflationary de Sitter points for the scalar potential. Both of them play a role of the uplifting term V_{uplift} in Eq. (2.5), and then the height of the inflationary de Sitter point $V_{\text{inf}} \sim H^2$ satisfies the relation:

$$\mathcal{O}(H^2) \sim \mathcal{O}(V_{\text{uplift}}) \sim \mathcal{O}(m_{3/2}^2). \quad (2.20)$$

This relation is similar to Eq. (2.8). Therefore, again we cannot combine the high scale inflation with the low scale SUSY breaking.

There are some models avoiding the relation (2.20). In Refs. [8] and [10], the Kähler potential of the modulus includes the α' -correction, and it changes the value of $\hat{\sigma}(f^I)$. These models can be solutions for a tension between the inflation scale and the SUSY breaking scale, however, both of them require the fine-tuning of parameters in the superpotential to separate the inflation scale from the SUSY breaking scale as is the case for the KL model. In Ref. [19], it was suggested that the SUSY breaking scale is much smaller than the inflation scale if the moduli roll into the large volume minimum after the inflation⁵. Although the way to separate the two scales is interesting, there is a overshooting problem because of the extreme difference between the inflation scale and the height of barrier. Therefore, the model requires to choose the initial condition precisely.

⁴One can find the derivation of this relation in Ref. [9].

⁵The Kähler potential in this model also contains the α' correction contribution, and then the inflationary de Sitter point is generated.

A simple solution for the overshooting problem is the positive exponent term discussed in Ref. [11]. However, such a positive exponent term prevent moduli from rolling into the large volume minimum, and then we cannot combine the positive exponent terms with the large volume models.⁶

3 Instant uplifted inflation

In this section, we propose a modulus inflation model which can separate the inflation scale from the SUSY breaking scale. In that model, there are two important ingredients. One is the existence of a scalar field which has the superpotential as follows:

$$W_{\text{uplifton}} = \mu_Y^2 Y e^{-c_Y T}. \quad (3.1)$$

In this paper, we refer to such a field Y as “uplifton”. As mentioned in section 2, the volume-type modulus inflation must add SUSY breaking terms for the realization of the inflationary de Sitter point. The F-term of the uplifton can be source of such a point, which decrease exponentially with the increasing VEV of modulus. This feature enables the separation of the two scales. The superpotential like Eq. (3.1) can arise, e.g., from the string instanton effects [20] or anomalous U(1) couplings [21]. Even in the effective theory of simple 5D SUGRA models on S^1/Z_2 , the factor $e^{-c_Y T}$ is always associated with bulk matter fields in the superpotential induced at one of the fixed point, if bulk matters are charged under the Z_2 odd U(1) gauge vectors [22].

We consider a model in which the superpotential has the form such as

$$W = \mu_Y^2 Y e^{-c_Y T} + A e^{-aT} - B e^{-bT}. \quad (3.2)$$

In this model, the F-term potential (2.3) is proportional to terms with negative exponents of T . Therefore, as the modulus rolls into the direction of the large VEV, the scale of the scalar potential decreases exponentially that makes the SUSY breaking scale today small enough irrespective of the magnitude of the initial inflationary scale. This mechanism seems similar to the one in Ref. [19] in which the modulus reach the large volume minimum. However, there is a broad distinction between them. In our model, the modulus doesn’t have to reach the extremely large VEV. As we will see in the following, the exponentially decreasing feature is important to combine this separation mechanism with the second important ingredient explained below.

The second key ingredient is the positive exponent term [11]. By virtue of the fact that the modulus doesn’t need to have an extremely large VEV, we can add the positive exponent terms in the superpotential⁷ such as

$$W_{\text{positive}} = \tilde{A} e^{aT}, \quad (3.3)$$

⁶We would like to thank Tetsutaro Higaki for pointing out this issue.

⁷Some moduli stabilization models with positive exponent terms are considered in Ref. [11] and their application to the inflation model can be found in Refs. [10],[11] and [13].

where

$$\tilde{A} = Ae^{-a'\langle S \rangle}.$$

The field S is a heavy modulus which is already stabilized at a higher scale than the cut off scale in our discussion. Such a positive exponent term can be generated, e.g., if we consider the gauge kinetic function which has a form:

$$f = w_S S + w_T T.$$

Then, if the gaugino condensation occurs for the $SU(N)$ gauge group with the above gauge kinetic function, the following superpotential term is generated:

$$W = Ae^{-\frac{2\pi}{N}(w_S S + w_T T)}. \quad (3.4)$$

As mentioned in Ref. [11], the coefficient w_T may have a negative value in some cases (e.g., in heterotic M theory [23], or magnetized D9-brane [24]). Such gauge kinetic functions also can be realized in 5D SUGRA models where the moduli mixing in gauge kinetic functions are determined by the arbitrary coefficients of cubic polynomials governing the structure of the $\mathcal{N} = 2$ gauge vector multiplets whose fifth components correspond to moduli S and T [25].

The uplifton and the positive exponent term can realize the separation of the two scales. We show the scalar potential which includes the two ingredients in Fig. 2 schematically. To make the mechanisms clear, let's consider a model in which the superpotential is given by

$$W = W_{\text{inf}} + W_{\text{min}}, \quad (3.5)$$

$$W_{\text{inf}} = Ae^{-aT} - Be^{-bT} + \mu_Y^2 Y e^{-c_Y T}, \quad (3.6)$$

$$W_{\text{min}} = w_0 - \tilde{C}e^{cT} + \mu_X^2 X. \quad (3.7)$$

We assume following conditions:

$$c_X T_{\text{inf}} \sim a T_{\text{inf}} \sim b T_{\text{inf}} \sim \mathcal{O}(1), \quad (3.8)$$

$$c T_{\text{min}} \sim \mathcal{O}(1), \quad (3.9)$$

$$T_{\text{min}} \sim 10 T_{\text{inf}}, \quad (3.10)$$

$$|\tilde{C}e^{cT_{\text{inf}}}| < |w_0| \ll |Ae^{-aT_{\text{inf}}}|, |Be^{-bT_{\text{inf}}}|, |\mu_Y^2 e^{-c_Y T_{\text{inf}}}|, \quad (3.11)$$

where T_{inf} denotes the typical VEV of T around the inflationary point, and T_{min} denotes the one around the minimum. Then, the scalar potential around the inflationary de Sitter point is dominated by W_{inf} . We can represent the scalar potential around the inflationary de Sitter point by

$$V|_{T \sim T_{\text{inf}}} \sim \frac{1}{8\sigma^3} (\mu_Y^4 e^{-2c_Y \sigma} + D_T \hat{W}_{\text{inf}} K^{T\bar{T}} D_{\bar{T}} \hat{W}_{\text{inf}} - 3|\hat{W}_{\text{inf}}|^2), \quad (3.12)$$

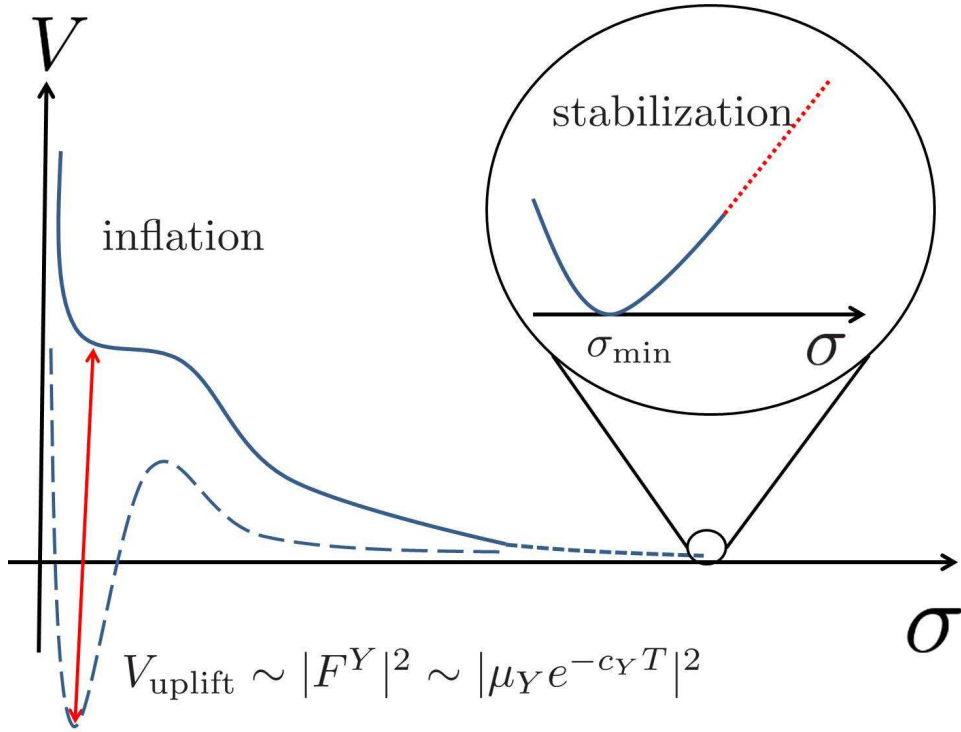


Figure 2: An schematic illustration of the instant uplifted inflation scenario explained in Sec. 3. The F-term of the uplifton may yield the inflationary de Sitter point at the high scale. The minimum is separated from that point, and the scale near the minimum is much smaller than the inflationary region. The potential includes the positive exponent term, then it will blow up for σ larger than σ_{\min} . The blowing up feature is drawn by the red dotted line.

where $\hat{W} \equiv W|_{X=Y=0}$.

The dominating part of the superpotential terms will change however, from W_{inf} to W_{min} after the moduli rolling down into the minimum. That is because the condition (3.10) leads $c_X T_{\text{min}} \sim a T_{\text{min}} \sim b T_{\text{min}} \sim \mathcal{O}(10)$, then $|W_{\text{inf}}|$ becomes an exponentially suppressed value. Therefore, the scalar potential around the minimum is mainly determined by W_{min} . In this situation, the scalar potential is represented by

$$V|_{T \sim T_{\text{min}}} \sim \frac{1}{8\sigma^3}(\mu_X^4 + D_T \hat{W}_{\text{min}} K^{T\bar{T}} D_{\bar{T}} \hat{W}_{\text{min}} - 3|\hat{W}_{\text{min}}|^2). \quad (3.13)$$

We don't have to care about the overshooting, hence \hat{W}_{min} contains a positive exponent term in the scalar potential and the only required condition for the superpotential W_{min} is found that it contains at least a single positive exponent term. Therefore we can consider some models with different types of stabilization potential aside from the positive exponent terms. We will see some illustrative models in the next section.

4 Some illustrative models

We show three explicit models to realize the mechanism discussed in Sec. 3. Those three models are different with respect to the stabilization potential. In all models, we use the same Kähler potential K , and the superpotential terms W_{inf} dominating inflation as follows:

$$K = -3 \log(T + \bar{T}) + |X|^2 - \frac{1}{\Lambda^2} |X|^4 + |Y|^2 - \frac{1}{\Lambda'^2} |Y|^4, \quad (4.1)$$

$$W = W_{\text{inf}} + W_{\text{min}}, \quad (4.2)$$

$$W_{\text{inf}} = C e^{-cT} - D e^{-dT} + \mu_Y^2 Y e^{-c_Y T}. \quad (4.3)$$

It is difficult to solve the dynamics of multiple fields simultaneously. Here we choose a set of parameters such that the masses of the fields other than the inflaton ($\sigma = \text{Re}T$) are heavy, and then treat the following models as the single field inflation model. In this case, we can analyze the relevant part of the whole dynamics based on the following effective potential:

$$V_{\text{eff}} = \frac{1}{8\sigma^3}(\mu_Y^4 e^{-2c_Y \sigma} + \mu_X^4 + K^{T\bar{T}} D_T \hat{W} D_{\bar{T}} \hat{W} + -3|\hat{W}|^2), \quad (4.4)$$

where $\hat{W} \equiv W_{\text{inf}}|_{X=Y=0} + W_{\text{min}}|_{X=Y=0}$. The detailed discussions are given in the appendix A. We use this effective potential in the following analyses.

4.1 KKLT type

We consider the model which contains the following superpotential terms W_{min} governing the whole dynamics around the minimum:

$$W_{\text{min}} = w_0 + A e^{aT} + \mu_X^2 X.$$

Then we choose the following set of parameters:⁸

$$\begin{aligned}
C &= (0.9)^2 \times 3 \times 10^{-5}, & D &= (0.9)^2 \times 1 \times 10^{-5}, & c &= \frac{\pi}{15}, & d &= \frac{\pi}{25}, & c_Y &= \frac{\pi}{70}, \\
w_0 &= (0.9)^2 \times 2 \times 10^{-11}, & A &= (0.9)^2 \times 2 \times 10^{-19}, & a &= \frac{\pi}{60}, \\
\mu_Y &= (0.9) \times 1.562633 \times 10^{-3} & \mu_X &= (0.9) \times 6.18 \cdots \times 10^{-6},
\end{aligned} \tag{4.5}$$

and we choose the initial condition:

$$\sigma(0) = 19.2, \quad \sigma'(0) = 0.$$

The smallness of the parameters C and D can be naturally realized if the moduli mixing occurs as $C = C' e^{-c_S S}$, $D = D' e^{-d_S S}$ and S is stabilized at a high scale. Therefore we don't consider the "smallness" of parameters as a consequence of a fine-tuning. We set the parameter μ_X in such a way that the AdS minimum is uplifted to the Minkowski minimum. We just admit a fine-tuning of the parameters μ_X and μ_Y which originates from the cosmological constant problem and the generation of the inflationary inflection point⁹, respectively. The solution of these fine-tuning problems is beyond the scope of this paper. Aside from these deep problems, we don't need the fine-tuned parameters to separate the inflation scale and the SUSY breaking scale as we have shown in the previous section.

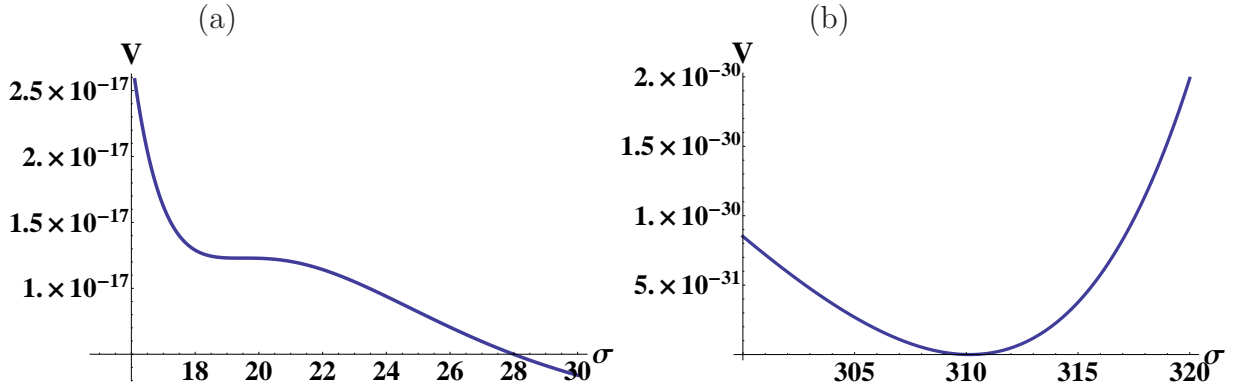


Figure 3: The scalar potential for the KKLT-type model (a) in the vicinity of the inflationary inflection point and (b) in the vicinity of the minimum.

We show the shape of the potential around T_{inf} and the one around T_{min} in Fig. 3. We can find that the height of the potential is extremely suppressed around the minimum

⁸The factors $(0.9)^n$ in Eq. (4.5) represent a rescaling to fit the WMAP normalization at the e-foldings $N \sim 50$ before the end of inflation.

⁹We can find the inflection point inflation in string theory e.g. [6], [10] and [26], and in MSSM inflation models [27].

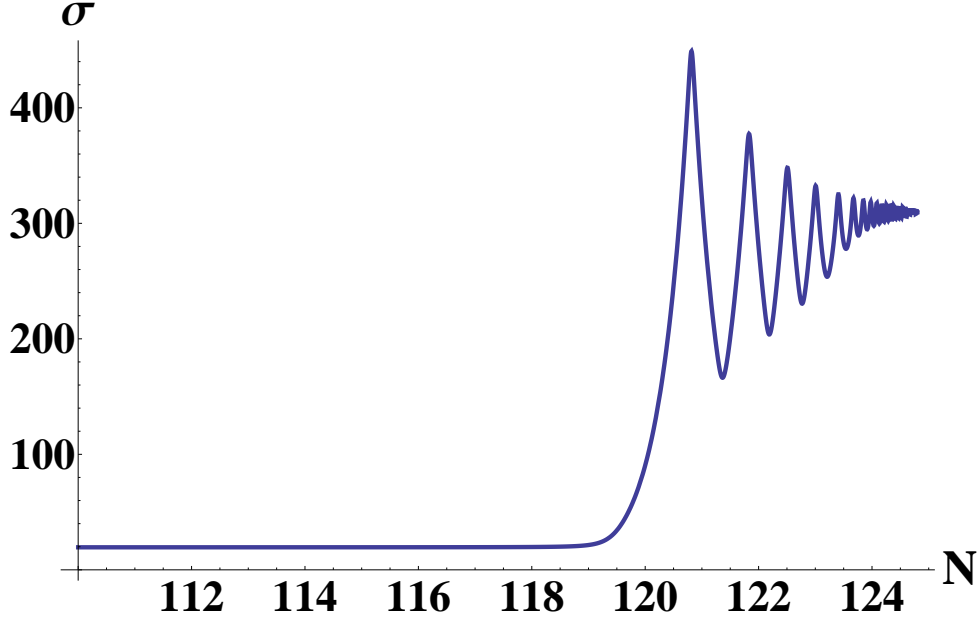


Figure 4: The evolution of the inflaton $\sigma = \text{Re}T$ for the KKLT-type model as functions of the e-folding number N in the last stage of the inflation.

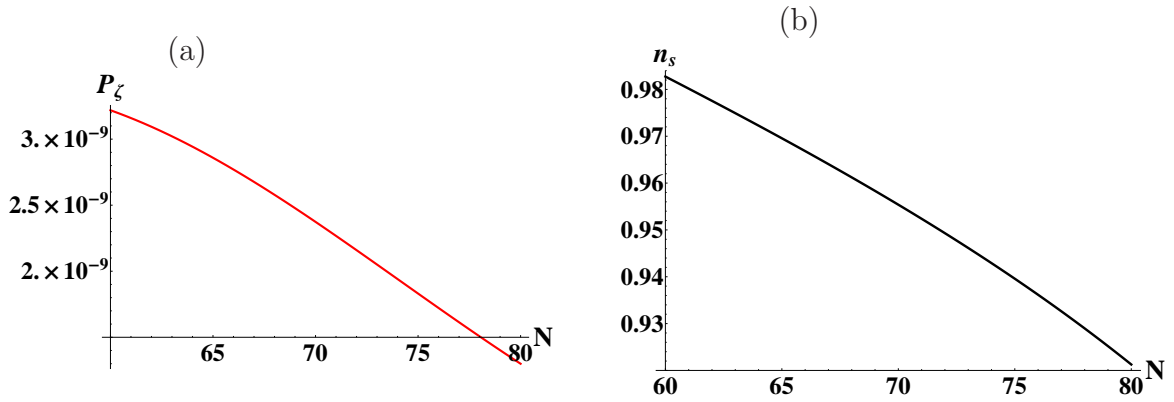


Figure 5: (a) The scalar power spectrum and (b) the spectral index of the scalar power spectrum for the KKLT type model as a function of the e-folding number. The region from 40 to 60 e-foldings before the end of the inflation is focused here.

compared with the one around the instantly uplifted de Sitter region $T \sim T_{\text{inf}}$. The evolution of the inflaton is shown in Fig. 4, and we find that the overshooting problem does not occur due to the positive exponent terms. The power spectrum of the scalar curvature perturbation $\mathcal{P}_\zeta = \frac{V}{24\pi^2\epsilon}$ and its spectral index $n_s = 1 + 2\eta - 6\epsilon$ in this model¹⁰ can be found in Fig. 5. The tensor to scalar ratio $r = 16\epsilon$ is $\mathcal{O}(10^{-10})$ in this model (and in the following two models discussed in Sec.4.2 and Sec.4.3), therefore these observables \mathcal{P}_ζ, n_s and r are consistent with the CMB observation [28].

We derive the SUSY breaking parameters around the minimum as follows:

$$m_{3/2} = 2.8 \text{ [TeV]}, \quad F^X = 4.8 \text{ [TeV]}, \quad \sqrt{K_{T\bar{T}}}F^T = 475 \text{ [GeV]}, \quad F^Y = 277 \text{ [GeV]}.$$

The gravitino mass is $\mathcal{O}(10^3 \text{ GeV})$ at the minimum, even though the Hubble parameter is $\mathcal{O}(10^9 \text{ GeV})$ during inflation. The hierarchy of $\mathcal{O}(10^6)$ is realized between the SUSY breaking scale and the inflationary one which confirms the success of the separation mechanism proposed in this paper.

4.2 Racetrack type

Next we adopt the racetrack-type superpotential terms

$$W_{\text{min}} = Ae^{aT} + Be^{-bT} + \mu_X^2 X. \quad (4.6)$$

We choose the set of parameters,

$$\begin{aligned} C &= (0.9)^2 \times 3 \times 10^{-5}, \quad D = (0.9)^2 \times 1 \times 10^{-5}, \quad c = \frac{\pi}{15}, \quad d = \frac{\pi}{25}, \quad c_Y = \frac{\pi}{70}, \\ A &= (0.9)^2 \times 5 \times 10^{-14}, \quad B = (0.9)^2 \times 6 \times 10^{-9}, \quad a = \frac{\pi}{175}, \quad b = \frac{\pi}{140}, \\ \mu_Y &= (0.9) \times 1.562633 \times 10^{-3}, \quad \mu_X = (0.9) \times 5.66 \cdots \times 10^{-6}, \end{aligned}$$

and the following initial conditions,

$$\sigma(0) = 19.4, \quad \sigma'(0) = 0.$$

These parameters in W_{inf} are almost the same as those in the subsection 4.1. So, the evolution of the modulus, the spectral index, the power spectrum, and the tensor to scalar ratio are similar to those in the previous subsection. The SUSY breaking parameters in this model is found as

$$\begin{aligned} m_{3/2} &= 2.4 \text{ [TeV]}, \quad F^X = 4.1 \text{ [TeV]}, \\ \sqrt{K_{T\bar{T}}}F^T &= 291 \text{ [GeV]}, \quad F^Y = 295 \text{ [GeV]}. \end{aligned}$$

Again the low scale SUSY breaking is realized and the separation is successful.

¹⁰Because this model is treated as the single field inflation with a non-canonical kinetic term generated by the Kähler potential (4.1), the slow roll parameters ϵ and η are given by $\epsilon = \frac{\sigma^2(\partial_\sigma V)^2}{3V^2}$, $\eta = \frac{2\sigma^2\partial_\sigma^2 V}{3V}$.

4.3 R-symmetric type

Finally, we consider the following superpotential terms:

$$W_{\min} = Ae^{aT} + \mu_X^2 X.$$

Then, we choose the following set of parameters,

$$\begin{aligned} C &= 3 \times 10^{-5}, & D &= 1 \times 10^{-5}, & c &= \frac{\pi}{15}, & d &= \frac{\pi}{25}, & c_Y &= \frac{\pi}{70}, \\ A &= 2 \times 10^{-12}, & a &= \frac{\pi}{370}, & \mu_Y &= 1.562651 \times 10^{-6}, & \mu_X &= 5.63 \cdots \times 10^{-6}, \end{aligned}$$

and the initial conditions,

$$\sigma(0) = 19.45, \quad \sigma'(0) = 0.$$

The observables during inflation are similar to the previous models, however, the SUSY breaking parameters differ substantially from the other ones. We show the SUSY breaking parameters in this model:

$$\begin{aligned} m_{3/2} &= 4.3 \text{ [TeV]}, & F^X &= 5.0 \text{ [TeV]}, \\ \sqrt{K_{T\bar{T}}} F^T &= 5.4 \text{ [TeV]}, & F^Y &= 430 \text{ [GeV]}. \end{aligned}$$

We find that the F-term of the modulus is comparable with that of the SUSY breaking sector X and the gravitino mass. The difference from the previous models is caused by the existence of R-symmetry. In this model, the superpotential W_{\min} has the exact R-symmetry, and the effective potential around the minimum has an approximate R-symmetry. In Ref. [29], it is pointed out that there is an R-symmetric SUSY breaking minimum for the modulus whose imaginary part is shifted under the R-symmetry transformation¹¹. Therefore, the modulus in this model also plays a SUSY breaking field, and the F-term of the modulus becomes relatively large. Such SUSY breaking parameters produce a different pattern of the superparticle spectrum. That is relevant to the particle phenomenology.

4.4 Comments on the difference between the instant uplifted inflation and the KL model

We give some comments on the difference between the instant uplifted inflation and the KL model. The KL model is attractive from the viewpoint of the simplicity, and the phenomenological consequences of the KL stabilization is studied recently. As shown in the Ref. [16], the volume modulus gets a heavy mass, and then the F-term of the modulus

¹¹We would like to thank Hiroyuki Abe for noticing this point.

is much smaller than that of the original KKLT modulus, then the contribution from such a F-term is much smaller than the anomaly mediation. Therefore the gaugino masses are dominated by the anomaly mediation, and the gravitino mass $m_{3/2}$ have to be $\mathcal{O}(100)\text{TeV}$ to get the sufficient gluino mass. On the other hand, the sfermions get the mass of order $m_{3/2}$ because they are not sequestered to avoid the tachyonic mass problem.

In our model, the modulus can get the relatively large F-term, and such a F-term can play an important role from the phenomenological viewpoint. For example, in the original KKLT type models (shown in Sec. 4.1 or 4.2), the anomaly mediation and the modulus mediation can be comparable, and then we can realize the mirage mediation [30] which is an attractive scenario because it solves the little hierarchy problem elegantly [31]. In our model, the F-term of modulus is determined by the W_{\min} and then we can obtain various patterns of phenomenological consequences. This feature is an important difference between the KL model and the instant uplifted inflation.

5 Conclusion

We proposed a new class of a mechanism to separate the scale of inflation with moduli fields from the SUSY breaking scale in this paper. The two ingredients are required to achieve the separation. One is the existence of “uplift” which has the following form of the F-term $F^Y = \mu_Y^2 e^{-c_Y T}$, then the inflationary de Sitter point can be realized in the scalar potential. Because the F-term decreases exponentially as T increases, we could make the minimum where the scale of the scalar potential is extremely smaller than the one during inflation. The other ingredients is the positive exponent term in the superpotential like $\tilde{C}e^{cT}$, which prevent the overshooting after inflation. As we have shown, we don’t need the fine-tuned parameters to separate the two scales. Due to the separation, we could adopt some different patterns of stabilization potential after inflation. Therefore we can make some different phenomenological models. For example, if we adopt the KKLT type model as the stabilization term, we may realize the mirage mediation [30], [31] which solves the SUSY little hierarchy problem.

In this paper, we focused on the separation between the inflation scale and the SUSY breaking scale. In order to construct realistic models, we have to combine these models with the successful Big-Bang nucleosynthesis. In addition, the low scale SUSY breaking models predict SUSY particles with TeV scale masses, and they may be discovered at the LHC in the near future. So, it is also important to analyse the prediction of such SUSY models. We will investigate concrete phenomenological models combined with the instant uplifted inflation proposed in this paper as a future work [32].

Acknowledgments

The author especially would like to thank Hiroyuki Abe for the early collaboration, useful discussions, and reading the manuscript carefully, and Tetsutaro Higaki for many useful comments and discussions. He is also grateful to Hajime Otsuka, Keigo Sumita, Yoshiyuki Tatsuta for discussions and comments.

A Derivation of the effective potential

The Kähler potential and superpotential in our model are given as follows:

$$K = -3 \log(T + \bar{T}) + |X|^2 - \frac{|X|^4}{\Lambda^2} + |Y|^2 - \frac{|Y|^4}{\Lambda'^2}, \quad (\text{A.1})$$

$$W = \mu_X^2 X + \mu_Y^2 Y e^{-c_Y T} + \hat{W}, \quad (\text{A.2})$$

$$\hat{W} = W_{\text{inf}}|_{X=Y=0} + W_{\text{min}}|_{X=Y=0}, \quad (\text{A.3})$$

where $\Lambda, \Lambda' \ll 1$, W_{inf} is given by Eq. (4.3) and W_{min} takes some patterns given by Eq. (4.5), Eq. (4.6), and Eq. (4.7).

The F-term scalar potential is generically given by

$$V = e^K (D_I W K^{I\bar{J}} D_{\bar{J}} \bar{W} - 3|W|^2). \quad (\text{A.4})$$

Then, we expand the potential in powers of X and Y , up to the quadratic terms. The terms of the 0th, 1st, and 2nd order of X and Y are represented respectively as $V^{(0)}, V^{(1)}, V^{(2)}$:

$$\begin{aligned} V^{(0)} &= \frac{1}{8\sigma^3} (\mu_X^4 + \mu_Y^4 e^{-2c_Y \sigma} + D_T \hat{W} K^{T\bar{T}} D_{\bar{T}} \hat{W} - 3|\hat{W}|^2) \\ &\equiv \frac{1}{8\sigma^3} V_0, \end{aligned} \quad (\text{A.5})$$

$$\begin{aligned} V^{(1)} &= \frac{1}{8\sigma^3} (K_T K^{T\bar{T}} D_{\bar{T}} \hat{W} - 2\hat{W}) \mu_X^2 X \\ &\quad + \frac{1}{8\sigma^3} (K^{T\bar{T}} D_{\bar{T}} \hat{W} (K_T - c_Y) - 2\hat{W}) \mu_Y^2 e^{-c_Y T} Y + \text{h.c.}, \end{aligned} \quad (\text{A.6})$$

$$\begin{aligned} V^{(2)} &= \frac{1}{8\sigma^3} \left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + 2\mu_X^2 + |\hat{W}|^2 \right] |X|^2 \\ &\quad + \frac{1}{8\sigma^3} \left[V_0 + \frac{4\mu_Y^4 e^{-2c_Y \sigma}}{\Lambda'^2} + (K^{T\bar{T}} (K_T - c_Y) (K_{\bar{T}} - c_Y) - 1) \mu_Y^4 e^{-2c_Y \sigma} + |\hat{W}|^2 \right] |Y|^2 \\ &\quad + \frac{1}{4\sigma^3} (1 + c_Y \sigma) \mu_X^2 \mu_Y^2 e^{-c_Y \bar{T}} X \bar{Y} + \frac{1}{4\sigma^3} (1 + c_Y \sigma) \mu_X^2 \mu_Y^2 e^{-c_Y T} Y \bar{X}. \end{aligned} \quad (\text{A.7})$$

The extremum conditions in terms of \bar{X} and \bar{Y} are found respectively as

$$\left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + 2\mu_X^2 + |\hat{W}|^2 \right] X + 2(1 + c_Y\sigma)\mu_X^2\mu_Y^2e^{-c_Y T}Y = 2(\sigma D_T\hat{W} + \hat{W})\mu_X^2, \quad (\text{A.8})$$

$$\begin{aligned} & \left[V_0 + \frac{4\mu_Y^4e^{-2c_Y\sigma}}{\Lambda'^2} + (K^{T\bar{T}}(K_T - c_Y)(K_{\bar{T}} - c_Y) - 1)\mu_Y^4e^{-2c_Y\sigma} + |\hat{W}|^2 \right] \\ & + 2(1 + c_Y\sigma)\mu_X^2\mu_Y^2e^{-c_Y\bar{T}}X = (2\hat{W} - K^{T\bar{T}}D_{\bar{T}}\hat{W}(K_{\bar{T}} - c_Y))\mu_Y^2e^{-c_Y\bar{T}}. \end{aligned} \quad (\text{A.9})$$

For a notational convenience, we define the following quantities:

$$V_{Y\bar{Y}} \equiv \left[V_0 + \frac{4\mu_Y^4e^{-2c_Y\sigma}}{\Lambda'^2} + (K^{T\bar{T}}(K_T - c_Y)(K_{\bar{T}} - c_Y) - 1)\mu_Y^4e^{-2c_Y\sigma} + |\hat{W}|^2 \right], \quad (\text{A.10})$$

$$V_{X\bar{X}} \equiv \left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + 2\mu_X^2 + |\hat{W}|^2 \right], \quad (\text{A.11})$$

$$V_{X\bar{Y}} \equiv 2(1 + c_Y\sigma)\mu_X^2\mu_Y^2e^{-c_Y\bar{T}}, \quad (\text{A.12})$$

$$V_{Y\bar{X}} \equiv 2(1 + c_Y\sigma)\mu_X^2\mu_Y^2e^{-c_Y T}, \quad (\text{A.13})$$

$$V_{\bar{X}}|_0 \equiv 2(\sigma D_T\hat{W} + \hat{W})\mu_X^2, \quad (\text{A.14})$$

$$V_{\bar{Y}}|_0 \equiv (2\hat{W} - K^{T\bar{T}}D_{\bar{T}}\hat{W}(K_{\bar{T}} - c_Y))\mu_Y^2e^{-c_Y\bar{T}}. \quad (\text{A.15})$$

Using these notations, Eqs. (A.8), (A.9) are represented by

$$\begin{pmatrix} V_{X\bar{X}} & V_{Y\bar{X}} \\ V_{X\bar{Y}} & V_{Y\bar{Y}} \end{pmatrix} \begin{pmatrix} X \\ Y \end{pmatrix} = \begin{pmatrix} V_{\bar{X}}|_0 \\ V_{\bar{Y}}|_0 \end{pmatrix}, \quad (\text{A.16})$$

that can be rewritten as

$$\begin{pmatrix} X \\ Y \end{pmatrix} \sim \frac{1}{V_{X\bar{X}}V_{Y\bar{Y}}} \begin{pmatrix} V_{Y\bar{Y}} & -V_{Y\bar{X}} \\ -V_{X\bar{Y}} & V_{X\bar{X}} \end{pmatrix} \begin{pmatrix} V_{\bar{X}}|_0 \\ V_{\bar{Y}}|_0 \end{pmatrix} \quad (\text{A.17})$$

$$= \begin{pmatrix} \frac{V_{\bar{X}}|_0}{V_{X\bar{X}}} - \frac{V_{\bar{Y}}|_0 V_{Y\bar{X}}}{V_{X\bar{X}}V_{Y\bar{Y}}} \\ \frac{V_{\bar{Y}}|_0}{V_{Y\bar{Y}}} - \frac{V_{\bar{X}}|_0 V_{X\bar{Y}}}{V_{X\bar{X}}V_{Y\bar{Y}}} \end{pmatrix}. \quad (\text{A.18})$$

Then, we can evaluate the VEV of X as follows:

$$X \sim \frac{2\mu_X^2(\sigma D_T \hat{W} + \hat{W})}{\left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + 2\mu_X^2 + |\hat{W}|^2\right]} - \frac{2\mu_X^2\mu_Y^4 e^{-2c_Y\sigma}(1 + c_Y\sigma)\{2\hat{W} - K^{T\bar{T}}D_T\hat{W}(K_{\bar{T}} - c_Y)\}}{\left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + |\hat{W}|^2\right]\left[V_0 + \frac{4\mu_X^4}{\Lambda^2} e^{-2c_Y\sigma} + |\hat{W}|^2\right]} \quad (\text{A.19})$$

$$\leq \frac{2\mu_X^2(\sigma D_T \hat{W} + \hat{W})}{\left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + 2\mu_X^2 + |\hat{W}|^2\right]} - \Lambda^2 \frac{2\mu_X^2(1 + c_Y\sigma)\{2\hat{W} - K^{T\bar{T}}D_T\hat{W}(K_{\bar{T}} - c_Y)\}}{\left[V_0 + \frac{4\mu_X^4}{\Lambda^2} + |\hat{W}|^2\right]}. \quad (\text{A.20})$$

We take account the relation $\mathcal{O}(\sigma D_T \hat{W}) \sim \mathcal{O}(\hat{W}) \sim \mathcal{O}(\sqrt{V_0})$. In the case $\mathcal{O}(\sqrt{V_0}) \geq \mathcal{O}\left(\frac{\mu_X^2}{\Lambda}\right)$, we find

$$X \leq \mathcal{O}\left(\frac{\mu_X^2}{\sqrt{V_0}}\right) \leq \mathcal{O}(\Lambda) \ll 1. \quad (\text{A.21})$$

On the other hand, in the case $\mathcal{O}(\sqrt{V_0}) \leq \mathcal{O}\left(\frac{\mu_X^2}{\Lambda}\right)$, we find

$$X \leq \mathcal{O}\left(\frac{\mu_X^2 \sqrt{V_0}}{(\frac{\mu_X^2}{\Lambda})^2}\right) \leq \mathcal{O}(\Lambda) \ll 1. \quad (\text{A.22})$$

As a result, we can always derive the relation $X \leq \Lambda$. From the similar discussion, we can find the relation $Y \leq \Lambda'$. These relations show that we can neglect the VEVs X and Y , and fluctuations of these fields around the VEVs have a large mass during inflation. Therefore we neglect the small VEVs and the fluctuations of X and Y , and find the effective potential V_{eff} for T shown in Eq. (4.4).

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